

**THE PLANCKIAN CONSPIRACY: STRING THEORY AND
THE BLACK HOLE INFORMATION PARADOX**DAVID A. LOWE¹*Department of Physics
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It has been argued that the consistency of quantum theory with black hole physics requires nonlocality not present in ordinary effective field theory. We examine the extent to which such nonlocal effects show up in the perturbative S-matrix of string theory.

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1. Introduction

The problem of information loss during quantum black hole evaporation [1,2] provides us with an important paradox which is likely to point the way to a better understanding of quantum gravity. If one takes the point of view that quantum mechanics is not drastically modified, an evaporating black hole must either: (i) encode the information carried by matter falling through the event horizon in the outgoing Hawking radiation and at the same time leave a freely falling observer crossing the event horizon essentially undisturbed [3,4,5], or (ii) store the information in a long-lived or stable remnant [6]. The second possibility suffers from the serious drawback that such remnants are likely to be infinitely produced in ordinary low energy scattering experiments [7]. The studies in two dimensions of semiclassical dilaton-gravity models [8] provide strong evidence that the first possibility is not realized in the context of local field theoretic degrees of freedom coupled to gravity.

For a large mass black hole, all local invariants outside and in the vicinity of the horizon are small for a large fraction of the lifetime of the black hole. In this case one might expect to be able to set up a local low energy effective field theory on this slowly varying spacetime, on a family of suitably defined slices. Further details of this argument, which we refer to as the nice slice argument, may be found in [9]. The usual rules of local quantum field theory tell us that local operators inside commute with local operators outside the horizon. Since the fields inside the horizon are correlated with the infalling state, an observer outside will conclude that information has been lost.

This argument suggests that if information is to be encoded in Hawking radiation, some new kind of nonlocality not present in local effective field theory is required. At the same time, these new nonlocal effects should not spoil ordinary low energy physics when black holes are not present. In this paper we further explore the idea that string theory realizes this possibility. This point of view has been considered previously in [10,11].

Because strings are extended objects, they are sensitive to nonlocal geometric invariants. This provides the needed loophole in the nice slice argument. In the black hole problem there is a large nonlocal invariant – the relative boost between the infalling and outgoing frame in which Hawking radiation is measured. We will show this leads to nonlocal effects over macroscopic spacetime separations visible in string S-matrix elements. We argue that these effects undermine the usual nice slice argument and that a nonlocal effective field theory is necessary to correctly describe low energy physics on a family of nice slices.

Another approach to this problem would be to compute the commutator of string fields in some suitable string field theory. In [12,13], it was argued that the commutator of free string fields in light-cone gauge was essentially local. This is no longer the case in the interacting theory [14], where the so-called transverse spreading was studied. In [9], this result is generalized to string fields with nontrivial separation in the longitudinal plane. There the nonlocal effect found is much larger than that present in the S-matrix. However, then one is faced with the difficult question of the interpretation of these off-shell quantities. It is possible these large nonlocal effects are simply a gauge artifact. On the other hand, these off-shell effects in flat space could show up in on-shell physical amplitudes computed around a finite mass black hole type background. Regardless of which interpretation turns out to be correct, the nonlocal behavior present in the S-matrix should be thought of as a

lower bound on the degree of nonlocality present in string theory.²

The paper is organized as follows. In section 2, tree-level S-matrix elements corresponding to the interaction of Hawking type particles with a particle freely falling across an event horizon are considered. These amplitudes fall off over a characteristic length scale that becomes macroscopically large at sufficiently late times. In section 3, we consider some of the physical consequences of these amplitudes and argue that they support the notion that degrees of freedom on a stretched horizon (a timelike surface lying outside the true event horizon) retain information about the infalling state. Section 4 discusses higher genus corrections to tree-level results, and we conclude with some brief remarks in section 5.

2. String Spreading in the Regge Limit

In the following we will be interested in computing string S-matrix elements between a set of operators at low energy relative to an observer freely falling into a large mass M_{bh} black hole, and another set of operators at low energy relative to an observer who stays outside the black hole and measures the Hawking radiation at late times. We will refer to such operators as “nice” operators. The stringy effects of interest here survive as the mass of the black hole goes to infinity, at which point the region of spacetime near the horizon is simply flat space. It suffices, therefore, to compute the string amplitudes in flat space, folding in appropriate wavepackets to construct the operators of interest. We denote by a the relative boost factor between the two frames. For a Schwarzschild background this increases exponentially with Schwarzschild time t_S , $a = e^{t_S/4M_{bh}}$, becoming very large at late times. We use the null coordinates X^+ , X^- to parametrize the longitudinal plane of the spacetime. The horizon corresponds to the line $X^+ = 0$, with $X^+ < 0$ outside the horizon. The $D - 2$ transverse coordinates are denoted \vec{X} and the conventions for the metric are $ds^2 = 2dX^+dX^- - (d\vec{X})^2$.

With respect to these Minkowski coordinates, the center of mass energy squared s of these amplitudes is very large at late times (of order a), while the momentum transfer squared t is small. In this limit the Regge behavior of the string amplitudes is apparent. Regge-Gribov techniques [15] may be used to compute the string amplitudes. This procedure has been considered previously in [16]. The amplitudes are factorized onto fictitious Reggeon particles with angular momentum $J = \alpha(t) = 2 + \alpha't/2$. Vertices coupling these Reggeons to ordinary particles such as gravitons, may be extracted by factorizing tree-level amplitudes. These vertices may then be sewn together to yield amplitudes at arbitrary numbers of loops.

We will adopt a simpler method, approximating the string path integral by a saddle-point calculation. This gives the correct asymptotic dependence on s for the amplitudes of interest here, but in general will not yield the full t dependence. Fortunately, the full t dependence will not be important in the following. The saddle-point approximation to the

² This of course assumes that we are treating the string degrees of freedom as fundamental. It remains a logical possibility that string theory may be formulated in terms of an underlying local theory.

string path integral has been considered before by Gross and Mende [17], in the context of large s , fixed angle (i.e. fixed t/s) scattering. Taking the fixed t limit of their expressions yields the amplitudes of interest here. It may be shown that when the determinant factors from the integration over fluctuations in the moduli are included, with the appropriate measures, the large s asymptotics coincide with those obtained via Regge-Gribov methods, despite the fact that for fixed $t < 1/\alpha'$ the saddle-point approximation is no longer strictly valid.

The string path integral is

$$A = \int \mathcal{D}X \exp\left(-\frac{1}{4\pi\alpha'} \int \partial X \bar{\partial} X + i \sum_j p_j \cdot X(z_j)\right) . \quad (2.1)$$

The saddle point is dominated by the worldsheet

$$X_{cl}^\mu = i\alpha' \sum_j p_j^\mu \log |z - z_j| , \quad (2.2)$$

which leads to the amplitude

$$A \sim \exp(-\alpha' E) , \quad (2.3)$$

where E may be thought of as the electrostatic energy of charges p_j^μ on the worldsheet at positions z_j

$$E = \sum_{i < j} p_i \cdot p_j \log |z_i - z_j| . \quad (2.4)$$

This energy is then to be minimized by varying the independent moduli.

As an example, consider the four-graviton scattering amplitude at tree-level. The mass-shell condition leads to the relation $s + t + u = 0$, where s, t, u are the Mandelstam variables: $s = (p_1 + p_2)^2$, $t = (p_1 + p_3)^2$ and $u = (p_1 + p_4)^2$. The energy E is invariant under $SL(2, C)$ transformations at tree-level, so depends on only one complex modulus

$$x = \frac{(z_1 - z_3)(z_2 - z_4)}{(z_1 - z_2)(z_3 - z_4)} . \quad (2.5)$$

Equation (2.4) then is rewritten

$$E = \frac{1}{2} (t \log |x| + u \log |1 - x|) . \quad (2.6)$$

Minimizing with respect to x then gives the solution

$$x_{cl} = -t/s , \quad (2.7)$$

so that in the limit $t \ll s$, $E \sim -\frac{1}{2}t \log s$. Note that because $x_{cl} \ll 1$ in this limit, the amplitude factorizes. This kind of factorization between the infalling operators and the outgoing Hawking operators will be a property of the generic N particle amplitudes.

Including the determinant from fluctuations in x , with the appropriate measure, the final amplitude is

$$A(s, t) \sim s^{2+\alpha' t/2} . \quad (2.8)$$

This amplitude could not be obtained from a local field theory with a finite number of fields, since (2.8) violates polynomial boundedness. There one would obtain at tree-level an amplitude polynomial in the momenta, multiplied by factors containing poles coming from internal propagators. One may then view this amplitude (2.8) as a field theory amplitude multiplied by a form factor corresponding to the effective size of the strings,

$$A(s, t) \sim s^2 F_s^2(t) , \quad (2.9)$$

where the form factor is

$$F_s(t) = \exp\left(\frac{\alpha' t}{4} \log s\right) . \quad (2.10)$$

This tells us the effective size of the strings is given by the Lorentz invariant expression

$$(\delta X^\mu)^2 \sim \alpha' \log s . \quad (2.11)$$

An important point to note is that if we consider the four-point amplitude of a nice low energy pair of Hawking particles ($p^+ \sim 1/a\ell$, $p^- \sim a/\ell$) in one frame and a nice low energy pair of infalling operators ($p^+ \sim 1/\ell$, $p^- \sim 1/\ell$) in the other frame, kinematics restricts the momentum transfer to be purely transverse, in the large s limit. Here ℓ is defined to be some typical experimental scale. This means the amplitude is not sensitive to the longitudinal spreading indicated in (2.11), i.e. $\delta X^+ \delta X^- \sim \alpha' \log s$. This turns out to be a general feature of the Regge amplitudes of fixed numbers of nice particles in the large s limit.

One way to ensure that the longitudinal momentum transfer $q^+ q^-$ is comparable to \bar{q}^2 , is to consider amplitudes with large numbers ($O(a)$) of nice operators. However, the extra large number in the problem generally leads to the break down of the above approximations. Alternatively one can consider amplitudes involving states with $M^2 \sim a$. These states may be thought of as long strings stretched between the Hawking probe and the infalling operators. These amplitudes can be computed within the saddle-point approximation, and do exhibit the longitudinal spreading indicated by (2.11), as will be shown in the following.

To simplify matters further, we will replace the large M^2 state by two graviton fields with momenta that are large with respect to both the infalling and Hawking frames, and consider the five-graviton amplitude corresponding to the production of these gravitons via a pair of Hawking operators interacting with an infalling operator. These two graviton fields are not “nice” operators, according to our previous definition, but nevertheless they may be produced via the interaction of some set of nice fields. This corresponds to an experiment in which the Hawking observer conducts a scattering experiment at low energy in her frame which involves a high-energy interaction with the infalling operator, which produces this high-mass intermediate state. The high-mass state then decays into a pair of high-energy gravitons. Of course this is not the most likely final state, which is expected

to involve a large number of particles. We should not be surprised then if this amplitude is highly suppressed – for now we are only interested in whether the longitudinal string spreading suggested by (2.11) shows up.

The amplitude is obtained by minimizing (2.4) with respect to the two independent complex moduli. Let us denote the momenta of the Hawking particles by the subscripts h and h' , those of the infalling particle by i , and the momenta of the outgoing pair of gravitons j and j' . The products of momenta satisfy: $p_h \cdot p_{h'} < O(1)$, with all other products between different momenta of order $a \gg 1$. The equations to be solved are

$$\sum_{k \neq l} \frac{p_l \cdot p_k}{z_l - z_k} = 0, \quad (2.12)$$

with no sum on l . Up to $SL(2, C)$ transformations, given the above conditions on the products of momenta, the solution satisfies $|z_h - z_{h'}| \sim O(1)$, with all other $|\delta z|$ of order a . This indicates the Hawking operators factorize with the exchange of a Reggeon coupling them to the other operators. Computing the amplitude leads to

$$A(s, t, M^2, \phi) \sim s^{\alpha' t/2} e^{-\alpha' M^2 f(\phi)/2}, \quad (2.13)$$

neglecting additional analytic factors from the measure and fluctuations in the moduli. Here s is of order a , $t = (p_h + p_{h'})^2$, $M^2 = (p_j + p_{j'})^2$, and $f(\phi)$ is a function of the scattering angle ϕ defined by $\sin^2(\phi/2) = -2p_i \cdot p_j / M^2$ which has the form

$$f(\phi) = -(\sin^2(\phi/2) \log \sin^2(\phi/2) + \cos^2(\phi/2) \log \cos^2(\phi/2)). \quad (2.14)$$

The function $f(\phi)$ will generically be of order one.

Again we see the appearance in (2.13) of the $s^{\alpha' t/4}$ form factor which is present whenever Reggeon exchange dominates the amplitude. The longitudinal components of the momentum transfer satisfy $q^+ \sim 1/a\ell$, $q^- \sim a/\ell$ so that $q^+ q^- \sim O(\bar{q}^2)$. Therefore, the momentum transfer is no longer purely transverse as in the usual Regge amplitudes, and the amplitude displays the longitudinal spreading $\delta X^+ \delta X^- \sim \alpha' \log s$ implied by equation (2.11). This will be a general property of diffractive scattering amplitudes of sets of nice operators which lead to a long string (i.e. a state with $M^2 \sim a$) in the final state.

Because this amplitude is suppressed by the factor $\exp(-\alpha' M^2 f(\phi)/2)$, one might worry that typical string configurations are not spread in this manner. If one considers inclusive amplitudes in which the different final states are summed over, the exponentially large number of final states with center of mass energy M will at least partially compensate for this suppression factor. Nevertheless, the suppression factor is independent of t , so we may fix M^2 , ϕ and vary t independently. The length scale over which the amplitude falls off as we vary t gives a measure of the typical longitudinal component of the spreading. Of course, on the basis of (2.11), one could argue such spreading simply follows from Lorentz invariance. The purpose of the above calculation is to demonstrate such spreading is in fact manifest in S-matrix elements.

Let us now consider further the spacetime properties of the form factor (2.10). The Fourier transform of the form factor yields the density of the extended object

$$\rho(\delta x) = \int dq e^{iq \cdot \delta x} F_s(q^2). \quad (2.15)$$

Substituting (2.10) into (2.15), implies that the typical size of the string is (2.11). In light-cone gauge this yields the usual $\log s$ transverse spreading [10,18,14]. Using (2.15) to compute the longitudinal spreading leads to

$$\langle \delta X^- \rangle \sim \alpha' \langle q^- \rangle \log s, \quad \langle \delta X^+ \rangle \sim \alpha' \langle q^+ \rangle \log s, \quad (2.16)$$

where $\langle q \rangle$ is the typical momentum transfer of the string under consideration. This is similar to the expression derived in [10] but contains an additional subleading $\log s$ factor. This implies that a string stops Lorentz contracting as its apparent length reaches the string scale, and actually begins growing logarithmically with the boost. These expressions should be contrasted with the longitudinal spreading one finds in the off-shell commutator in light-cone gauge. The power law fall-off with X^- in that case implies no length scale is associated with the longitudinal spreading. See [9] for further details.

If one folds in wavepackets with a typical spread ℓ (in space and time, for example), and we work in the center of mass frame of the string, a stringy uncertainty relationship is obtained

$$\delta X^\mu \sim \frac{\alpha'}{\ell} \log s + \ell. \quad (2.17)$$

This implies there is a minimum observable length scale obtainable in performing fixed momentum transfer, high energy scattering experiments

$$\delta X_{min}^\mu \sim \sqrt{\alpha' \log s}. \quad (2.18)$$

Similar uncertainty relationships have been considered before in [16,17].

3. Physical Consequences

Consider an accelerating observer who hovers outside the horizon of a large mass M_{bh} black hole, such that she undergoes a proper acceleration $1/\rho_0$, and performs experiments on length scale ℓ . The path the observer follows is then

$$x^- = \rho_0 e^{t_S/4g^2 M_{bh}}, \quad x^+ = -\rho_0 e^{-t_S/4g^2 M_{bh}}, \quad (3.1)$$

where t_S is the time measured by the observer and g is the string coupling constant. We take the infalling operators to be localized near $x^- = 0$. As argued above, including the effects of the finite size of wavepackets, the infalling and Hawking operators will interact over a range

$$x^- < \alpha' \langle q^- \rangle \log a + 1/\langle q^+ \rangle = \frac{\alpha' t_S}{4g^2 M_{bh} \ell} e^{t_S/4g^2 M_{bh}} + \ell e^{t_S/4g^2 M_{bh}}, \quad (3.2)$$

or equivalently

$$\rho_0 < \frac{\alpha' t_S}{4g^2 M_{bh} \ell} + \ell. \quad (3.3)$$

The first term in (3.3) is the inherently stringy spreading, while the second term is just the usual field theoretic spreading, arising from the finite size of the detector. In principle, ℓ can be made as small as one wishes so that this field theoretic overlap may be neglected. However, as ℓ is made very small, the effective size of the string increases. One concludes the nonlocal terms are always significant when

$$\rho_0 < \sqrt{\frac{\alpha' t_S}{4g^2 M_{bh}}} . \quad (3.4)$$

If we choose to view string theory in terms of a set of degrees of freedom local with respect to some fixed target space background, the above result suggests that degrees of freedom on a stringy stretched horizon $x^+ x^- = \alpha' \log(a)$ carry information about the infalling state. The size of this horizon differs by a factor $\log(a)$ with the analogous result of [10].

Note the above calculations were based on perturbative string theory. As we will see later, in the fixed momentum transfer, large s limit strong coupling effects become important when $s > 1/(\alpha' g^2)$, where g is the string coupling constant. We conclude the above result (3.4) is only applicable for time

$$t < O(4M_{bh} g^2 \log 1/g^2). \quad (3.5)$$

A detector must either undergo string scale accelerations or cross over the path of the infalling body, to attempt to detect these degrees of freedom with this kind of experiment.³ This does not directly lead to any breakdown of low energy effective field theory, unless we are prepared to extrapolate the tree-level results beyond their range of validity (3.5). However, in the case of a finite mass black hole, a similar stretched horizon will be present, which couples to low energy observers via Hawking radiation. It seems plausible that these information carrying degrees of freedom on the stretched horizon will couple to the Hawking modes, allowing information to leak off the stretched horizon. In terms of a low energy effective field theory on a set of nice slices, these effects will show up as nonlocal interactions between the Hawking region and the infalling region.

On the other hand, since string theory appears to allow us to ask questions involving infalling and outgoing operators at the same time, if we assume the picture suggested by above is valid at very late times, there would appear to be a problem with large scale violations of the equivalence principle. Any observation on the Hawking radiation which yields information about the state fallen into the black hole, would seem to lead to a violent perturbation of the infalling observer as she passes through the horizon.

Of course, the fact that the amplitude (2.13) is highly suppressed means such processes occur extremely rarely. Consequently, these processes can transfer only a small amount of information to the outgoing Hawking radiation. However, we regard the existence of these rare information transferring perturbative processes as an important hint that string theory may realize possibility (i) of the Introduction.

³ This assumes $\log 1/g^2$ is roughly of order 1. If $\log 1/g^2 \gg 1$ then the string spreading effects come into play at substring-scale accelerations.

Clearly something more subtle must happen if information retrieval before the end-point of black hole evaporation is to be implemented in string theory. It is possible the new physics enters in at the nonperturbative level. The perturbative picture described above, in which stringy degrees of freedom carry information on the stretched horizon about the infallen state, will only provide a correct description of the physics for sufficiently early times (3.5). After that time, much larger effects must set in if all the information of the infalling state is to eventually reach an outside observer in the form of Hawking radiation. To be consistent with known physics, the infalling observer cannot perceive these effects until she approaches the singularity. This suggests that a better understanding of the stringy physics near the singularity is required to determine whether all information escapes encoded in the Hawking radiation.

4. Extension to Higher Genus

At arbitrary genus G , a saddle-point analysis similar to the above gives an estimate of the large s asymptotics of the fixed t amplitude. Let us restrict attention to the genus G four-graviton amplitude. This is dominated by a term which comes from sewing together tree-level amplitudes

$$A_G(s, t) \sim \prod_{i=1}^{G+1} A_{\text{tree}}(s, t_i) , \quad (4.1)$$

which is maximized subject to the condition that $\sum \sqrt{-t_i} \geq \sqrt{-t}$. It follows from the form of the tree-level amplitudes that the maximum occurs when the momentum transfer is equally shared between the different internal legs, so that $\sqrt{-t_i} \sim \sqrt{-t}/(G+1)$. When the determinants from the fluctuations of the moduli with the correct measures are included [17], one finds

$$A_G(s, t) \sim (g^2 s)^{G+1} i^G (G+1)^{9G} s^{\alpha' \frac{t}{2(G+1)} + 1} . \quad (4.2)$$

Neglecting the effect of the $(G+1)^{9G}$ factor for the moment, we see that $g^2 s$ is an effective expansion parameter in the large s , fixed t regime, and we should expect strong coupling effects to become important when $\alpha' s > 1/g^2$.

Taking into account the $(G+1)^{9G}$ behavior, the sum of the leading terms

$$A_{\text{sum}}(s, t) \sim \sum_{G=0}^{\infty} A_G(s, t) , \quad (4.3)$$

will diverge for fixed s , t and g . However, it may be Borel resummed along the lines of [19], to yield a convergent expression with the same asymptotic expansion in g as (4.3)

$$A_{\text{resum}}(s, t) = \int_0^\infty dz e^{-z} \sum_{G=0}^{\infty} \frac{z^{9N}}{(9N)!} A_G(s, t) . \quad (4.4)$$

We may evaluate the Borel resummation (4.4) using a saddle-point approximation [19] which yields

$$|A_{\text{resum}}(s, t)| \sim e^{-\sqrt{-2\alpha' t \log s \log(1/g^2)}} . \quad (4.5)$$

The value of G at the saddle point is

$$G_0 + 1 = \sqrt{\frac{-\alpha' t \log s}{2 \log(1/g^2)}} , \quad (4.6)$$

which must be large ($G_0 \gg 1$) for the saddle-point approximation to be valid. Equation (4.5) then is valid when

$$\frac{-2}{\alpha' t} \log(1/g^2) \ll \log s . \quad (4.7)$$

If $\frac{-2}{\alpha' t} \log(1/g^2) \gg \log s$, the tree-level term will dominate the series, which leads to the $(\delta X)^2 \sim \alpha' \log s$ spreading described in the previous sections.

Since we have summed only the leading terms, one might worry that subleading terms might change the asymptotic behavior given by (4.5). Subleading corrections may be estimated and resummed using the above procedure as described in [19]. The dominant corrections appear to come from integrals over the moduli and take the form $G^3/t A_G(s, t)$. In the resummed amplitude, these will be negligible if $G_0^3/t \ll 1$ which yields the additional constraint $|\alpha' t/2|^{1/3} \log s \ll \log(1/g^2)$. Therefore the above approximation to the resummed amplitude (4.5) is at best only valid in the range

$$\left(\frac{1}{g^2}\right)^{2/\alpha'|t|} \ll \alpha' s \ll \left(\frac{1}{g^2}\right)^{(2/\alpha'|t|)^{1/3}} . \quad (4.8)$$

This region in parameter space will only exist when $|t| > \alpha'/2$, so this is telling us about the dense central region of string that appears as the outer extremities undergo the logarithmic spreading described by the tree-level term [18]. Of course, it has not been proven that the subleading terms are negligible when all subleading corrections are taken into account. It remains a possibility that these terms conspire to eliminate the range of validity of (4.5).

For now, let us assume the resummed amplitude is valid for some range with $\log s > \frac{2}{\alpha'|t|} \log(1/g^2)$. Equation (4.5) may then be interpreted as a product of form factors, as in (2.9), now with higher genus corrections included. The resummed form factor falls off more slowly for large $|t|$, than the tree-level result. This slower falloff turns out to be consistent with the bound of Cerulus and Martin [20] on the fastest possible falloff in a local field theory, indicating that the resummed amplitude behaves in a more local way than just the tree-level term alone. Consequently, strings are more dense in the central region than would be expected from the tree-level result. Arguments presented in [21] suggest this behavior does not persist for $\alpha' s > 1/g^2$. There it was argued a much faster falloff must appear if string theory is to be consistent with the Bekenstein bound on the maximum amount of information within a fixed volume.

5. Conclusions

In this paper, we have studied certain perturbative string S-matrix elements relevant to the black hole information problem. These matrix elements correspond to rare events which transfer information from an infalling body to the Hawking radiation in a nonlocal

manner, disturbing the infalling body in the process. Such processes do not appear in a local effective field theory which might be expected to describe low energy physics on a family of nice slices of a large mass black hole. We therefore regard the existence of such processes in string theory as evidence that the effective field theory must contain nonlocal interactions. This nonlocality undermines the usual argument for information loss.

The perturbative results we have found are only valid for sufficiently early times, so cannot be the whole story. It is possible nonperturbative effects play an important role at later times. These effects must be much larger than the perturbative effects described above, if all infalling information is to escape in the form of Hawking radiation.

We have argued that string theory supports the idea that degrees of freedom on a stretched horizon retain information about what fell into the black hole. These degrees of freedom may then be interpreted as a kind of stringy hair. This is allowed, because when fields with spin greater than 2 are present, the usual no hair theorem of classical general relativity breaks down. Such hair is usually associated with some conserved quantity. It would be interesting to identify the symmetries of string theory which lead to these conserved quantities.

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References

- [1] S. W. Hawking, *Phys. Rev.* **D14** (1976) 2460.
- [2] S. W. Hawking, *Comm. Math. Phys.* **43** (1975) 199.
- [3] L. Susskind, L. Thorlacius, and J. Uglum, *Phys. Rev.* **D48** (1993) 3743.
- [4] G. 't Hooft, *Nucl. Phys.* **B335** (1990) 138; *Phys. Scr.* **T36** (1991) 247, and references therein.
- [5] E. Verlinde and H. Verlinde, *Nucl. Phys.* **B406** (1993) 43; K. Schoutens, E. Verlinde, and H. Verlinde *Phys. Rev.* **D48** (1993) 2690; Y. Kiem, H. Verlinde and E. Verlinde, hep-th/9502074.
- [6] Y. Aharonov, A. Casher and S. Nussinov, *Phys. Lett.* **191B** (1987) 51.
- [7] S. B. Giddings, *Why aren't black holes infinitely produced?*, hep-th/9412159, and references therein.
- [8] C. G. Callan, Jr., S. B. Giddings, J. A. Harvey and A. Strominger, *Phys. Rev.* **D45**, (1992) R1005; S. W. Hawking, *Phys. Rev. Lett.* **69** (1992) 406; B. Birnir, S. B. Giddings, J. A. Harvey and A. Strominger, *Phys. Rev.* **D46** (1992) 638; T. Banks, A. Dabholkar, M.R. Douglas and M. O' Loughlin, *Phys. Rev.* **D45** (1992) 3607; J.G. Russo, L. Susskind and L. Thorlacius, *Phys. Lett.* **B292** (1992) 13, *Phys. Rev.* **D46** (1992) 3444, *Phys. Rev.* **D47** (1993) 533; A. Bilal and C. G. Callan, *Nucl. Phys.* **B394** (1993) 73; S. P. de Alwis, *Phys. Lett.* **289** (1992) 278; D. A. Lowe, *Phys. Rev.* **D47**, (1993) 2446.
- [9] D. A. Lowe, J. Polchinski, L. Susskind, L. Thorlacius and J. R. Uglum, to appear.
- [10] L. Susskind, *Phys. Rev. Lett.* **71** (1993) 2367; *Phys. Rev.* **D49** (1994) 6606.
- [11] L. Susskind and L. Thorlacius, *Phys. Rev.* **D49** (1994) 966.
- [12] E. Martinec *Class. Quan. Grav.* **10** (1993) L187; *Strings and Causality*, preprint EFI-93-65, hep-th/9311129.
- [13] D. A. Lowe, *Phys. Lett.* **B326** (1994) 223.
- [14] D. A. Lowe, L. Susskind, and J. Uglum, *Phys. Lett.* **B327** (1994) 226.
- [15] V.N. Gribov, *Sov. Phys. JETP* **26** (1968) 414.
- [16] D. Amati, M. Ciafaloni and G. Veneziano, *Phys. Lett.* **197B** (1987) 81; *Int. Jour. Mod. Phys.* **A3** (1988) 1615; *Phys. Lett.* **216B** (1989) 41.
- [17] D.J. Gross and P.F. Mende, *Phys. Lett.* **197B** (1987) 129; *Nucl. Phys.* **B303** (1988) 407.
- [18] M. Karliner, I. Klebanov and L. Susskind, *Int. Jour. Mod. Phys.* **A3** (1988) 1981; A. Mezhlumian, A. Peet and L. Thorlacius, *Phys. Rev.* **D50** (1994) 2725.
- [19] P. Mende and H. Ooguri, *Nucl. Phys.* **B339** (1990) 641.
- [20] F. Cerulus and A. Martin, *Phys. Lett.* **8** (1963) 80; A. Martin, *Nuovo Cimento* **37** (1965) 671.
- [21] L. Susskind, *The World as a Hologram*, Stanford University preprint SU-ITP-94-33, September 1994, hep-th/9409089.